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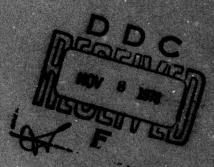
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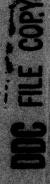
University of Maryland
College Park, Maryland 20742

and

S.L. Ossakow

Geophysical and Plasma Dynamics Branch
Plasma Physics Division

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We present a two-dimensional nonlinear stabilization mechanism for the EXB instability occurring in striated F region ionospheric plasma clouds. A simple model shows that the saturated power spectrum for striations $\propto k^{-2}$. Comparisons with experimental data from DNA sponsored STRESS barium cloud experiments and the results of numerical simulations are given.

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NONLINEAR STABILIZATION OF THE EXB GRADIENT
DRIFT INSTABILITY IN IONOSPHERIC PLASMA CLOUDS

The artificial injection of ion clouds in the ionosphere has been used as a technique to study various ambient conditions in the medium; e.g., the mapping of electric fields based upon their bulk motion in the medium (Haerendel and Lust, 1968). These clouds are found to display a typical pattern of development soon after their release. As they EXB drift in the crossed ambient electric and magnetic fields, they steepen on the backside (with respect to this motion), elongate and then striate on the steepened backside (Rosenberg, 1971). These striations are highly field-aligned and assume the shape of protruding fingers in their late time history, extending from the backside to the frontside. The motion of the ion cloud has been studied in detail by Haerendel et al. (1967). The steepening aspect has been explained by Perkins et al. (1973). The striations which result from the steepening have been generally attributed to the EXB gradient drift instability (Linson and Workman, 1970; Völk and Haerendel, 1971) originally applied by Simon (1963) to laboratory gas discharges. The numerical simulation studies performed by various groups to simulate the interaction of the plasma cloud and the ionosphere have agreed with the dynamics of the observed gross motion of the cloud and the development of fine structure on its steepened backside (Zabusky et al., 1973; Lloyd and Haerendel; 1973, Goldman et al., 1974; Scannapieco et al., 1976; Doles et al., 1976). In the work of Scannapieco et al. (1976), it was found that the striations, due to EXB instability obeyed a power law for the Fourier wavenumbers going as $\sim k^{-2}$. The experimental observations and numerical simulation studies indicate towards a coherent

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development of these striations in their nonlinear state (fingers, etc.). The recent DNA sponsored STRESS barium cloud experiments have provided data that agrees with the above interpretation and numerical simulation results. (Kelley et al., 1978; Baker and Ulwick, 1978).

In the present paper, we wish to present a theory for the coherent nonlinear development of the EXB instability in plasma clouds. In this theory the instability generates linearly damped harmonics through non-linear interaction and stabilizes. In fact, it has been shown sometime back that the so-called cross-field instability in the equatorial E region electrojet evolves nonlinearly through this stabilization mechanism (Rognlien and Weinstock, 1974). The difference in the EXB instability in the electrojet and in the clouds released at high F region altitudes is that in the electrojet region the differential motion between the Hall-drifting electrons and collisional non-magnetized ions results in the instability while in the high altitude case the difference arises through the large Pedersen mobility of ions as compared to that of electrons. Likewise, a gravity can also provide a differential motion between the species by providing ions a relative motion over electrons and would result in the collisional Rayleigh-Taylor instability which can be saturated by similar effects (Chaturvedi and Ossakow, 1977). The similarity between the EXB gradient drift instability in ionospheric plasma clouds and the collisional Rayleigh-Taylor instability in equatorial Spread F was noted in the numerical simulation work (Scannapieco and Ossakow, 1976; Scannapieco et al., 1976) and in the experimental observations (Kelley et al., 1978).

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The nonlinear set of equations describing the physical interaction of a plasma cloud in the ionosphere has been discussed many times in the past (Volk and Haerendel, 1971; Perkins et al., 1973; Scannapieco et al., 1976). We shall consider the simple case of a one dimensional plasma cloud. We further consider a local region on the steepened side of the cloud and study what happens to the nonlinear evolution of the perturbations, linearly unstable in this region due to the EXB instability. It should be pointed out here that this approach puts some restrictions on the validity of the results vis-a-vis real situations. First, a two-dimensional cloud of finite extent is not stationary and is continuously getting distorted as it moves in the ionosphere. However, on the faster time scale on which the striations grow, an assumption of one dimensionality gives us a quasi-stationary solution for the cloud. Second, we ignore the interaction of the cloud with the background ionosphere (the second level) which is a good enough assumption for the large clouds (one level model). In our geometry the earth's magnetic field, $\underline{B}_0$, is in the z-direction and the electric field is pointing along the x-axis. The steepening of the cloud in this configuration is such that for the backside the gradient is negative, and on the frontside it is positive (along the y-direction). The coupled set of nonlinear equations can be written as (see Perkins et at., 1973)

$$\frac{\partial n}{\partial t} - \frac{c}{B_0} \nabla_\perp n \cdot \nabla_\perp \bar{\phi} \times \hat{e}_z - \frac{2cT}{eB_0 \kappa_{en}} \nabla_\perp^2 n = 0 \tag{1}$$

and

$$\nabla_\perp \cdot (n \nabla_\perp \bar{\phi}) + \frac{2T}{e} \nabla_\perp^2 n = 0 \tag{2}$$

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where $\kappa_{\alpha n} = \Omega_\alpha / \nu_{\alpha n}$, $\Omega_\alpha = |e_\alpha| B_0 / m_\alpha c$, $\perp$ denotes perpendicular to $\underline{B}_0$, and $\hat{e}_z = \underline{B}_0 / |\underline{B}_0|$. Here eqn. (1) is the ion continuity equation in which the second term represents the effects due to $\underline{E} \times \underline{B}$ convection of plasma along a density gradient and the last term represents the ambipolar diffusion damping. (Here we have not taken into account the electron-ion Coulomb collisions which restrict the applicability of the results to lower F region altitudes.) The perturbations are assumed electrostatic and the generalized potential $\bar{\phi}$ is given by

$$\bar{\phi} = \phi - \frac{T}{e} \frac{1-\Delta}{1+\Delta} \ln\left(\frac{n}{n_0}\right) \tag{3}$$

where T is the temperature (in energy units) of the species (equal for ions and electrons), n_0 is a representative equilibrium plasma density and Δ is a small quantity : $\Delta = \frac{\kappa_{in}}{\kappa_{en}} = \frac{m_e}{m_i} \frac{\nu_{en}}{\nu_{in}} \ll 1$. The electrostatic potential ϕ includes an ambient part ϕ_0 and the perturbed part $\tilde{\phi}$:

$\phi = \phi_0 + \tilde{\phi}$, so

$$\nabla_\perp \phi = -\underline{E}_0 + \nabla_\perp \tilde{\phi}$$

Equation (2), a statement of quasi-neutrality, $\nabla_\perp \cdot \underline{J}_\perp = 0$, is obtained by subtracting the ion continuity equation from the electron one, and resembles the Poisson's equations. If written in terms of undisturbed and perturbed quantities, one can rewrite eqn. (1) and (2) as

$$\frac{\partial \tilde{n}}{\partial t} - \frac{c}{B_0} \nabla_\perp \tilde{\phi} \times \hat{e}_z \cdot \nabla_\perp n_0 - \frac{2cT}{eB_0 \kappa_{en}} \nabla_\perp^2 n = \frac{c}{B_0} \nabla_\perp \tilde{\phi} \times \hat{e}_z \cdot \nabla_\perp \tilde{n} \tag{4}$$

and

$$\nabla_\perp \cdot (n \nabla_\perp \tilde{\phi}) + \frac{T}{e} \nabla_\perp^2 n = \underline{E}_0 \cdot \nabla_\perp n \tag{5}$$

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where we have transformed to a frame moving with $c\underline{E}_o \times \underline{B}_o / B_o^2$ velocity. In equation (4), the second term gives the linear growth of perturbations and the third one represents diffusion damping. The term on the right hand side is a nonlinear term representing a nonlinear flux of particles through mode-interactions. Similarly in the potential equation (5), the left hand term has a linear and a nonlinear term. The linear term simply gives the perturbed potential, as it arises due to the differential mobilities between ion and electrons across a density perturbation (due to finite Pedersen effects on ions) in an ambient electric field, $\underline{E}_o$. The second term would be a nonlinear term. For one-dimensional perturbations, it is the dominant nonlinearity. In this case, the important nonlinear effect is the shielding of the driving electric field. It should be remembered here that the eqn. (5) is a lowest order equation in $\left(\dfrac{\nu_{in}}{\Omega_i}\right)$ and thus these quantities do not appear explicitly in the equation. A comparison between the two nonlinear terms in equations (4) and (5), readily shows that the term in (5) would be down by $\left(\dfrac{\nu_{in}}{\Omega_i}\right)$ as compared to the term in eqn. (4). This is of course valid only for two-dimensional perturbations, since for one dimensional perturbations, the term in equation (4) is negligibly small. But the problem is clearly inherently two-dimensional, and thus we adopt this approach here. (For more discussion on this point, see Rognlien and Weinstock, 1974; Chaturvedi and Ossakow, 1977.) If one linearizes eqn. (4) and (5) and Fourier analyses the perturbations as $\sim \exp(i\underline{k}\cdot\underline{x} - i\omega t)$ one obtains a familiar growth of perturbations at a rate given by

$$\gamma = \frac{cE_o}{B_o |L|} i \left(\frac{k_x}{k_\perp}\right)^2 - \frac{\nu_{en}}{\Omega_e \Omega_i} c_s^2 k_\perp^2 \tag{6}$$

where $\qquad L = \dfrac{n_o}{(dn_o/dy)}$ and $c_s^2 = 2T/m_i$

One can therefore rewrite eqns. (4) and (5) as

$$\frac{\partial \tilde{n}}{\partial t} = \gamma \tilde{n} + \frac{c}{B_o} \nabla_\perp \tilde{\phi} \times \hat{e}_z \cdot \nabla_\perp \tilde{n} \tag{4a}$$

and

$$n_o \nabla_\perp^2 \tilde{\phi} = \underline{E}_o \cdot \nabla_\perp n - \frac{T}{e} \nabla_\perp^2 n \tag{5a}$$

Obviously, eqn. (5a) has been written in its linearized form.* If one substitutes a two-dimensional perturbation in the nonlinear term of (4a), of the form: $\dfrac{\tilde{n}}{n_o} = A_{1,1} \sin(k_x x - \omega t) \cos k_y y$; one sees that the result is $\sim \left[\dfrac{cE_o}{2B_o} \dfrac{k_x^2}{k_\perp^2} k_y A_{1,1}^2\right] \sin 2k_y y$. That is, a second spatial harmonic in the y-direction (along the density gradient) is generated. We represent it as: $A_{2,o} \sin 2k_y y$. Thus our general perturbation has the form

$$\frac{\tilde{n}}{n_o} = A_{1,1} \sin(k_x x - \omega t) \cos k_y y + A_{2,o} \sin 2k_y y \tag{7}$$

From (5a), the associated potential perturbation is

$$\tilde{\phi} = i\beta \frac{\tilde{n}}{n_o}, \qquad \beta = \frac{k_x E_o}{k_\perp^2} \tag{8}$$

A substitution of (7) and (8) in (4a) leads to the coupled equations in mode amplitudes:

$$\frac{\partial A_{1,1}}{\partial t} = \gamma_{1,1} A_{1,1} - \alpha A_{1,1} A_{2,o} \tag{9}$$

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*For long wavelengths, of the order of hundreds of meters, the second term on right can be neglected. For $E_0 \approx 10mV/m$, $T = 10^3 \, {}^\circ K$, this term assumes importance at wavelengths of the order of tens of meters. In any case the term would cause a shift in the real part of ω.

and

$$\frac{\partial A_{2,0}}{\partial t} = \gamma_{2,0} A_{2,0} + \frac{\alpha}{2} A_{1,1}^2 \qquad (10)$$

With the coupling coefficient given by $\alpha = \frac{k_x^2 k_y}{k_\perp^2} \frac{cE_0}{B_0}$. If the mode $A_{2,0}$ is damped linearly, then $\gamma_{2,0}$ would be negative. Thus the nonlinear interactions cascade energy from the end where it is initially input (the linearly growing mode) to the k-space where linear theory predicts stability. This introduces a nonlinear damping for the linearly driven modes (see eqn. (9)), and one can think of a stationary state when the linear growth is balanced by the nonlinear damping. In this situation $\partial A_{1,1}/\partial t = \partial A_{2,0}/\partial t = 0$ and one obtains

$$A_{2,0} = \left| \frac{\gamma_{1,1}}{\alpha} \right| \sim \frac{1}{k_y |L|} \qquad (11)$$

and

$$A_{1,1} = \left(+ \frac{2|\gamma_{2,0}|}{\alpha} A_{2,0} \right)^{\frac{1}{2}}$$

$$\sim \frac{(2|\gamma_{1,1}||\gamma_{2,0}|)^{\frac{1}{2}}}{\alpha} \qquad (12)$$

Note that if $|\gamma_{1,1}| > |\gamma_{2,0}|$, one obtains that $A_{2,0} > A_{1,1}$. This means that if the spatial harmonics along the density gradient are relatively weakly damped, they are nonlinearly driven to larger amplitudes compared to the linearly driven modes. Further, these spatial harmonics

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can be thought to be quasi-linearly modifying the equilibrium density gradient and thus stabilizing the instability. Note that the saturated amplitude of the $A_{2,o}$ mode obtained here is of the same form as that obtained in the case of the collisional Rayleigh-Taylor instability (Chaturvedi and Ossakow, 1977). This indicates a close similarity between the driving effects for the collisional Rayleigh-Taylor instability in the equatorial F region and the EXB instability in F region Ba clouds. The numerical simulation of the nonlinear phase of the two instabilities also shows similar characteristics (Scannapieco and Ossakow, 1976) which is not surprising noting the similar form of the two coupled equations in the two cases. Experimental data of Kelley et al. (1978) also agrees with this interpretation. Substitution of typical numbers into (11) shows that for wavenumbers such that for $k_y L \sim 8$

$$A_{2,o} \simeq 12\% \tag{13}$$

We note here that the experimental observations by Kelley et al. (1978) and Baker and Ulwick (1978) are in line with the above estimate. Also, the recent one level numerical calculations give a saturation amplitude of this order [B. E. McDonald 1978, private communication.]

Further, it may be noted that the nonlinear generation of harmonics with large amplitudes increases the harmonic content in the final density profile and would steepen it (as has been found in the numerical calculations of Rognlien and Weinstock (1974) for their

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many-mode case) and the power law for the mode numbers goes as $\sim k^{-2}$ (eqn. (11)). This is also in agreement with the numerical simulations of <u>Scannapieco et al.</u> (1976) and the experimental results of <u>Kelley et al.</u> (1978). Similar observations of k^{-2} spectra for coherent equatorial F region irregularities (<u>Dyson et al.</u>, 1974) can also be interpreted in terms of a similar nonlinear theory (<u>Chaturvedi and Ossakow</u>, 1977). Finally, we remark that the next related problem to investigate would involve a two-level case (see <u>Scannapieco et al.</u>, 1976), in which image striations in the background ionosphere and the associated nonlinear effects are taken into account.

Acknowledgement

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DR. E. BAUER

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UNIVERSITY OF PITTSBURGH
PITTSBURGH, PA. 15213

DR. N. ZABUSKY
DR. M. BIONDI